

# Constraining interacting dark energy models with flux destabilization

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## Abstract

A destabilization in the transfer energy flux from the vacuum to radiation, for two vacuum decay laws relevant to the dark energy problem, is analyzed using the Landau-Lifshitz fluctuation hydrodynamic theory. Assuming thermal (or near thermal) equilibrium between the vacuum and radiation, at the earliest epoch of the Universe expansion, we show that the law due to renormalization-group running of the cosmological constant term, with parameters chosen not to spoil the primordial nucleosynthesis scenario, does soon drive the flux to fluctuate beyond its statistical average value thereby distorting the cosmic background radiation spectrum beyond observational limits. While the law coming from the saturated holographic dark energy does not lead the flux to wildly fluctuate, a more realistic non-saturated form shows again such anomalous behavior.

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Oftentimes, the present state of accelerated expansion of the Universe is related to some mysterious dark energy sector. This is linked to the longstanding cosmological-constant (CC) problem [1], by adding to it two distinct (but connected) difficulties: (i) the “new” CC problem, a puzzle of why the CC is small but non-zero, and (ii) the “cosmic coincidence”/“why now?” problem [2], a puzzle concerning the current near coincidence of the CC energy density,  $\rho_\Lambda$ , with that of matter despite they scale at different rates with expansion. It seems today that a new aspect in dealing with the CC problems lies in the landscape of string theory [3], though making predictions in such a theory constitutes an enormous challenge [4].

Long before the “environmental” variable CC approach of the string theory landscape it was noticed that, up to some extent, a traditional running of the CC can ameliorate the fine-tuning problems inherent to the CC by providing a viable mechanism to efficiently relax it from a very high value at the early Universe to its current tiny value. It was subsequently noted that some of the running CC models could successfully mimic the popular quintessence models as well as shed some light on the coincidence problem, thus becoming viable cosmological models of dark energy of the Universe. Arguably, the most appealing amongst them are those whose laws for the CC running can be inferred from some underlying physical theory. So, some of the most attractive dark energy models involve the CC running laws derived from quantum theory of particle fields on the classical gravitational background [5], quantum gravity [6] and gravitational holography [7]. A comprehensive list of phenomenological laws for the CC variation under consideration well before the discovery of dark energy, can be found in [8].

Barring a time-dependent gravitational coupling [6, 9, 10] or going over some scalar-tensor theory [11], the simplest way to achieve the infrared (IR) screening of the CC is through its decay into matter and/or radiation. The interesting possibility in this context, put forward long ago, was to consider the cosmological vacuum decay into radiation as a measure of the temperature of the vacuum [12]. Thus, if the interactions between the two components at the earliest moments of the Universe expansion were fast enough to bring them in thermal or near thermal equilibrium, then both -the vacuum and radiation- would share a common temperature for, at least some time, during the expansion. Interestingly, such a scenario would preclude a thermal equilibrium between the vacuum and the event horizon (which naturally occurs in these scenarios for most of the history of the Universe), for much of

the time except near the Planck time. Indeed, since the temperature of the horizon as given by the Hubble parameter  $H$ , and the temperature of the vacuum/radiation  $T$ , scale differently with the expansion, they may therefore coincide (during a radiation dominated epoch),  $H \simeq T$ , only when  $T \simeq M_{Pl}$ .

At a macroscopic level, the transfer of energy from vacuum to radiation (and vice versa) is governed by the continuity equation

$$\dot{\rho}_\Lambda + \dot{\rho}_R + 4H\rho_R = 0 , \quad (1)$$

with  $\rho_R$  being the radiation energy. Once the equilibrium (or near equilibrium) between the sub-systems vacuum/radiation gets established it will remain so provided the heat capacity of the whole system is positive-definite. Since the radiation heat capacity is necessarily positive, this amounts to having the heat capacity of vacuum

$$C_\Lambda = T \left( \frac{\partial S_\Lambda}{\partial T} \right)_V \quad (2)$$

positive, where  $S_\Lambda$  represents the entropy of a variable CC. (Admittedly, there is some ambiguity when taking the partial derivative in Eq. (2) as the volume should be kept constant but the latter depends on  $T$  in an expanding Universe).

Our target in the present paper is to study the fluctuations of the flux  $\dot{\rho}_\Lambda$  entering Eq. (1) around its statistical average value, from the laws emerging from the RG-running and gravitational holography, taking the macroscopic criterion that  $C_\Lambda > 0$  as a consistency condition. To fulfill this aim we shall employ the well known Landau-Lifshitz (LL) fluctuation hydrodynamic theory [13], which applies to equilibrium and nonequilibrium classical statistical theory [14]. Our particular emphasis will be on finding a scale dependence of these fluctuations, with a stabilization criterion that the root mean square of the fluctuations must never exceed the average value of the corresponding flux.

According to Landau and Lifshitz, if the flux  $\dot{y}_i$  of a given thermodynamic quantity, which evolves in a generic dissipative process, is governed by  $\dot{y}_i = \sum_j \Gamma_{ij} Y_j + \delta\dot{y}_i$  and the entropy rate obeys  $\dot{S} = \sum_i (\pm Y_i \dot{y}_i)$ , then the second moments of the fluctuations of the fluxes are given by  $\langle \delta\dot{y}_i \delta\dot{y}_j \rangle = (\Gamma_{ij} + \Gamma_{ji}) \delta_{ij} \delta(t_i - t_j)$ . The angular brackets denote statistical

average with respect to the reference state (i.e.,  $\Sigma_j \Gamma_{ij} Y_j$  which is supposed to be steady or quasi-steady and constitutes the systematic part of the flux), and the fluctuations  $\delta \dot{y}_i$  are considered spontaneous departures from that state, thus  $\langle \delta \dot{y}_i \rangle$  vanishes identically. The quantities  $\Gamma_{ij}$  and  $Y_i$  stand for the phenomenological transport coefficients and the thermodynamic force conjugate to the flux  $\dot{y}_i$ , respectively. In the expression for  $\dot{S}$  the minus sign must be taken when the product  $Y_i \dot{y}_i$  is negative, otherwise the plus sign should be considered. This theory has been successfully employed to constrain models for the decay of the cosmological constant into radiation and/or matter [15] as well as in the analysis of second order nonequilibrium phase transitions in isolated black holes [16].

The LL theory when applied to the decay of a variable CC includes just a single flux,  $\dot{\rho}_\Lambda(t)$ , governed by <sup>1</sup>

$$\dot{\rho}_\Lambda = \Gamma Y + \delta \dot{\rho}_\Lambda , \quad (3)$$

where the fluctuations,  $\delta \dot{\rho}_\Lambda$ , coming from the decay of the vacuum are assumed to be *Gaussian*, random fluctuations, with uncorrelated Fourier modes due to statistical homogeneity and isotropy. The thermodynamic force conjugate to the flux  $\dot{\rho}_\Lambda$  follows from combining the entropy production rate

$$\dot{S}_\Lambda = Y \dot{\rho}_\Lambda , \quad (4)$$

with Eq. (3). Finally, the second moment (i.e., the root mean square) of the fluctuations of the flux is given by

$$\langle \delta \dot{\rho}_\Lambda(t_i) \delta \dot{\rho}_\Lambda(t_j) \rangle = 2 \Gamma \delta(t_i - t_j) . \quad (5)$$

For the models explored below we have that

$$\dot{\rho}_\Lambda \sim \rho_\Lambda H . \quad (6)$$

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<sup>1</sup> The second law of thermodynamics implies the presence of particle creation and therefore the only flux to be considered is  $\dot{\rho}_\Lambda$ .

Near  $t = t_{Pl}$ , we are in the realm of  $a \sim t^\alpha$  cosmologies (see below) thereby both terms on the right hand side of Eq.(6) vary less and less faster with time. In addition, for the RG-running model, both  $\rho_\Lambda$  and  $H$  will approach a constant value at late times but well before the present moment (see below). Thus, the approximate steady-state regime can be maintained even for  $t \rightarrow \infty$ , provided  $C_\Lambda > 0$ .

Using the above sketched LL theory for the fluctuations of the fluxes, we aim to study their behavior for the approach to the CC problem based on the RG [5]. The RG is a conventional theoretical tool for investigating quantum effects and the scale dependence of a certain quantity. From the viewpoint of quantum theory of matter fields in curved space [17, 18], the renormalizability of the theory forces the vacuum action to contain the CC term as well as fourth derivative terms. Then, the CC term is viewed as a parameter subject to RG running and therefore it is expected to run with the RG scale, usually identified with an expansion quantity evolving smoothly enough to comply with the cosmological data. In such theories, therefore, even a “true” CC cannot be set to any definite fixed value (including zero) owing to the RG running effects. It may be surprising, however, that the time-dependence of the CC may be due to quantum effects from the RG, considering the familiar quadratic decoupling of heavy matter fields at low energy.<sup>2</sup> The reason for this result [20] lies in the high dimensionality ( $mass^4$ ) of the scaling quantity  $\rho_\Lambda$ , with the outcome that the more massive a field is, the more dominant the role it plays in the running -irrespective of the scale. Consequently, the running becomes stronger than logarithmic, thus providing a dynamical, efficient relaxation mechanism for the CC to go down the tiny value we observe today. Further, the above scenario for the CC running, with the choice for the RG scale  $\mu = H$ , may also furnish a viable cosmological model of dark energy [21, 22]. The strongest phenomenological constraints on the RG model (within a framework where the running of the CC goes at the expense of the energy transfer between vacuum and dark matter) have been obtained recently by analyzing density perturbations for the running CC [23] and considering the validity of the generalized second law for the running CC scenario [24].

The solutions for the RG-running vacuum decay into radiation for flat space read [21, 22],

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<sup>2</sup> Strictly speaking, the quadratic form of decoupling can be proved in a rigorous way only for higher derivative terms of the vacuum action, but not for the CC term itself [19]. Yet, usually, the same is assumed for the CC term.

$$\rho_R = \rho_{R0} a^{-4(1-\nu)} , \quad (7)$$

$$\rho_\Lambda = \left( \rho_{\Lambda0} - \frac{\nu}{1-\nu} \rho_{R0} \right) + \rho_{R0} \frac{\nu}{1-\nu} a^{-4(1-\nu)} , \quad (8)$$

$$H^2 = \frac{8\pi}{3} M_{Pl}^{-2} \left[ \left( \rho_{\Lambda0} - \frac{\nu}{1-\nu} \rho_{R0} \right) + \rho_{R0} \frac{1}{1-\nu} a^{-4(1-\nu)} \right] , \quad (9)$$

where  $\nu = \frac{\sigma}{12\pi} \frac{M^2}{M_{Pl}^2}$  is a dimensionless mass parameter driving the RG running. Here  $M$  represents an additive mass contribution of all virtual massive particles,  $\sigma = \pm 1$  depending on whether the highest-mass particle is a boson or a fermion, and the zero subscript denotes present-day value. Accordingly,  $|\nu| \sim 10^{-2}$  would signal the existence of a particle with Planck mass (or the existence of somewhat less massive particles with large multiplicities);  $|\nu| \sim 1$  would indicate the existence of a particle with trans-planckian mass;  $|\nu| \sim 10^{-6}$  would mean the existence of a particle with mass at the GUT-scale, whereas much smaller values of  $|\nu|$  would imply an approximate cancellation between bosonic and fermionic degrees of freedom.

For adiabatic vacuum decays one expects the Stefan law to be preserved (in the following we shall omit numerical factors since it is only the functional dependence that matters). Using Eq. (7) and bearing in mind that  $\rho_R \propto T^4$ , we get

$$T \sim a^{-1+\nu} , \quad (10)$$

and since  $T$  is not expected to increase with expansion we infer that  $0 < \nu < 1$ .

The entropy production associated to the CC decay is given by Gibb's equation with the chemical potential set to zero

$$\dot{S}_\Lambda = \frac{1}{T} V \dot{\rho}_\Lambda , \quad (11)$$

where  $V \sim a^3$ . From Eqs. (11) and (4) we obtain

$$Y = \frac{V}{T} \sim a^{4-\nu} . \quad (12)$$

A combination of  $\dot{\rho}_\Lambda \sim a^{-4(1-\nu)}H$ , from Eq. (8), with Eq. (3) yields

$$\Gamma \sim a^{5\nu-8}H . \quad (13)$$

We then determine the dimensionless ratio between the second moment of the fluctuations and the (squared) flux

$$\eta \equiv \frac{\langle \delta\dot{\rho}_\Lambda \delta\dot{\rho}_\Lambda \rangle}{\dot{\rho}_\Lambda \dot{\rho}_\Lambda} \sim \frac{\Gamma}{(\dot{\rho}_\Lambda)^2} = a^{-3\nu}H^{-1} . \quad (14)$$

Notice that the scaling dependence of the ratio  $\eta$  is crucial. Should  $\eta$  increase with expansion, sooner or later the fluctuations of the flux would become larger than the statistical average value of the flux, signaling destabilization -i.e., the flux would exhibit an erratic, unphysical, behavior. This would mean that there is not longer guarantee that the fluctuations of the flux preserve the equilibrium relations of the adiabatic decay of the vacuum, especially the Stefan law [25]. Further, this would seriously upset the black-body spectrum of the cosmic microwave background (CMB) beyond the limits allowed by observation. The opposite instance ( $\eta$  decreasing with expansion), corresponds to the usual, stable condition.

Using  $H \sim a^{-2(1-\nu)}$  from Eq. (9) for  $a \ll 1$ , we get  $\eta \sim a^{-5\nu+2}$ . Then, by requiring that  $\eta$  does not increase with expansion, we obtain

$$\nu \geq \frac{2}{5} . \quad (15)$$

On the other hand, for  $a \gg 1$  we have that  $H \sim \text{constant}$  whereby  $\nu > 0$  follows. From Eqs. (2) and (11) we get

$$C_\Lambda \sim \frac{4\nu}{1-\nu} a^{3\nu} , \quad (16)$$

whence for the allowed range of values for  $\nu$  the heat capacity of the vacuum results positive-definite and thermal equilibrium between vacuum and radiation is to hold for ever.

Before proceeding, it is expedient to check whether the systematic part of the flux is quasi-steady. It will be whenever the time scale for the vacuum to decay into radiation is larger than the expansion time (i.e.,  $\rho_\Lambda / |\dot{\rho}_\Lambda| \gtrsim H^{-1}$ ). Noting that, as follows from Eq. (8), the systematic part is  $\dot{\rho}_\Lambda = -4 \rho_{R0} \nu a^{-4(1-\nu)} H$ , it is readily seen this is guaranteed for  $\nu \gtrsim 3/4$ .

Next, we compare our LL bound, given by (15), with the existing bounds on the RG model at any epoch. As is observationally known, at the primordial nucleosynthesis time the ratio  $\rho_\Lambda/\rho_R$  did not exceed 0.05 [26]. Since

$$\frac{\rho_\Lambda}{\rho_R} \simeq \frac{\nu}{1-\nu}, \quad (17)$$

as follows from Eqs. (7) and (8) for  $a \ll 1$ , this sharply contrasts with the LL bound (15).

The above formulae also apply in the case of a dynamical CC scenario generically dubbed “holographic dark energy” (HDE). Originally derived for zero-point energies [7] as a bound on  $\rho_\Lambda$ , the saturated form of the HDE is usually written as [27]

$$\rho_\Lambda = \frac{3 M_{Pl}^2}{8\pi} c^2 L^{-2}, \quad (18)$$

where  $L$  denotes the size of the region (providing an IR cutoff) and  $c^2$  is a dimensionless constant. This is a very important concept since for  $c^2$  values of the order of unity, the HDE model also provides a very elegant solution of the “old” CC problem. Thanks to the relationship between the ultraviolet (UV),  $\rho_\Lambda \sim \Lambda^4$ , and the IR cutoff, the holographic information is consistently encoded in the conventional quantum field theory. The choice  $L = H^{-1}$  is clearly the most natural and simple possibility [28, 29, 30]. Then, with the aid of Friedman’s equation we can write

$$c^2 = \frac{1}{1+r_0}, \quad (19)$$

where  $r_0 = \rho_{R0}/\rho_{\Lambda0}$ .



In order to connect the above formulas for the RG case to Eq. (18), we shall write  $\rho_\Lambda$  from Eq. (8) in a different fashion, namely,

$$\rho_\Lambda = C_0 + C_2 H^2 \quad (20)$$

with

$$C_0 = \rho_{\Lambda_0} - \frac{3\nu}{8\pi} M_{Pl}^2 H_0^2, \quad C_2 = C_0 + \frac{3\nu}{8\pi} M_{Pl}^2. \quad (21)$$

The system of equations (20)-(21) is thus equivalent to the (7)-(9) set. The HDE law (18), follows from (20)-(21) by setting  $C_0 = 0$ . Then, Eqs. (18) and (19) are readily recovered -modulo, the obvious identification  $\nu = c^2$ . Now, as seen from Eq. (19), the bound from the LL theory,  $c^2 \geq \frac{2}{5}$ , is respected since the ratio  $r_0$  is tiny today,  $\sim 10^{-5}$ . Note that although the HDE law, Eq. (18) with (19), describing the vacuum decay, does intrinsically satisfy the LL bound, it does disturb the big bang nucleosynthesis scenario by a wide margin. The main problem is that the ratio  $\rho_R/\rho_\Lambda$  stays frozen during the whole cosmic expansion so that, for parameters driving the accelerated expansion of the Universe at late times, a transition to a radiation-dominated Universe is not feasible. Thus, the model cannot be considered realistic.

A more realistic class of models, which do allow transitions between the cosmological eras, is provided by the non-saturated HDE concept [29, 30, 31]. The parameterization is again given by Eqs. (18) and (19), but with  $c^2$  promoted to a function of cosmic time whence the ratio between the energy densities becomes a function of time. A criterion for a realistic non-saturated HDE model is to saturate the holographic bound asymptotically,  $c^2(t \rightarrow \infty) = 1$ , while having  $c^2 < 1$  in the radiation-dominated era. This type of parametrization of  $\rho_\Lambda$  in a non-saturation regime is particularly appealing since it reduces directly to (18), where again only the genuine IR cutoff shows up. It is easy to find a particular model that does not comply with the LL bound and the big bang nucleosynthesis bound simultaneously. It is possible to find such a function  $c^2$ , which does satisfy the above criterion for a realistic non-saturated HDE model, so that the RG law (20)-(21) become equivalent to the law obtained from the non-saturated HDE. Indeed, the choice

$$r(t) = r_0 \frac{a^{-4+\epsilon}}{1 - \alpha r_0 (1 - a^{-4+\epsilon}) (1 - \alpha + r_0)^{-1}} \quad (22)$$

with  $\epsilon = 4\alpha/(1 + r_0)$  and  $\alpha \equiv C_2 H_0^2 / \rho_{\Lambda_0}$  reproduces a non-saturated HDE model

$$\rho_\Lambda = \frac{3 M_{Pl}^2}{8\pi} c^2(t) L^{-2} , \quad (23)$$

with  $c^2 = \frac{1}{1+r(t)}$ , equivalent to the RG-running law given by Eqs. (20)-(21). As shown above, this model does not satisfy the LL bound without seriously affecting the big bang nucleosynthesis scenario.

In summary, two of the most popular and viable dark-energy models, based on vacuum decaying laws, appear largely compromised by flux destabilization. Our analysis relies on the assumption that, at some time close to the Planck era, the vacuum was in thermal (or near thermal) equilibrium with radiation. We have shown that the said equilibrium would persist at the time of big bang nucleosynthesis, where the vacuum is restricted to a tiny fraction of the total energy density [26]. If, in the meantime, the fluctuations of the energy flux become erratic (i.e.,  $\eta > 1$ ), then the radiation component will no longer present a black-body spectrum and the CMB will get seriously distorted. For a running CC scenario it has been shown that it is not possible to simultaneously hinder the growth of fluctuations (relative to the average value) and reduce the vacuum contribution at nucleosynthesis' time to an acceptable level. For a saturated HDE model, the flux remain under control at all times. However, if the nucleosynthesis bound is fulfilled, then an accelerated expansion at late times cannot be achieved. The non-saturated HDE model shows identical anomaly as the RG model.

We may conclude by saying that either a thermal equilibrium between vacuum and radiation did never occur or the dark energy models here considered are in need of revision.

## Acknowledgments

This work was partly supported by the Ministry of Science, Education and Sport of the Republic of Croatia under contract No. 098-0982887-2872, the Spanish Ministry of

Education and Science under Grant FIS 2006-12296-C02-01, and the “Direcció General de Recerca de Catalunya” under Grant 2005 SGR 00087.

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